JHU-TIPAC 98013

Strongly Interacting Neutrinos and the Highest Energy Cosmic Rays

G. Domokos and S. Kovesi-Domokos
Department of Physics and Astronomy
The Johns Hopkins University
Baltimore, MD 21218*

December 1998

Abstract

Cosmic rays of energies larger than the Greisen–Zatsepin–Kuzmin (GZK) cutoff may be neutrinos if they acquire strong interactions due to a "precocious unification" of forces. A scenario for this to happen is outlined. There is no contradiction with precision measurements carried out at LEP and SLAC. Observable consequences at LHC and future neutrino detectors are discussed.

PACS 13.85.T, 13.15, 14.60.S

A substantial number of cosmic ray events has been detected in which the primary energy appears to exceed the Greisen-Zatsepin-Kuzmin (GZK) cutoff, see *e.g.* [1, 2, 3, 4] and Szabelski's recent review [5]. The onset of the GZK cutoff itself is somewhat uncertain: the energy of the rapid turnover of the spectrum depends on several details (for instance, on the injection spectrum, *etc.*), see ref. [6] for a modern treatment. Nevertheless, it is unlikely that all the events reported can be explained by fluctuations in the shower development.

*E-mail: skd@haar.pha.jhu.edu

A number of explanations of the "anomalous" events has been offered; for an overview, *cf.* the proceedings of the 1997 University of Maryland workshop [7]. In a recent article, however, Burdman, Halzen and Gandhi ref. [8] conclude that none of the explanations offered in the literature is a convincing one.

Given this situation, we reexamine a proposal put forward some time ago, [9, 10]. In those papers, we proposed that at sufficiently high energies, neutrinos (in general, leptons) acquire some unspecified strong interaction and cause (possibly) post-GZK showers. Since neutrinos are neutral, have small magnetic moments and at low energies they have nothing but weak interactions, they can freely propagate through the 2.7°K background. Consequently, they can reach us from cosmological distances.

There were two weaknesses inherent in this suggestion. First, no mechanism has been offered as to how the neutrinos get their strong interaction. Second, due to the fact that we used a sharp (Θ function) threshold to turn on the strong interaction, the resulting amplitude violated unitarity, *cf.* [8].

The scenario we present here is based on recent work designed to overcome the hierarchy problem. The authors of refs. [11, 12, 13, 14, 15, 16] conjecture that unification may take place at a much lower energy than originally suspected, perhaps at a few TeV (*"precocious unification"*). Even though there are some problems with this approach (*e.g.* the long lifetime of the proton is still lacking a convincing explanation), it offers some interesting possibilities regarding the nature of the physics beyond the Standard Model.

The high unification masses obtained within the Standard Model and its minimal supersymmetric extensions, see [17, 18, 19, 20], is the result of the assumption of a "desert" between – approximately – the weak scale and the GUT scale. If one wants to accomplish unification at a lower energy, one has to postulate the existence of a rapidly increasing density of states somewhere above the weak scale.

For the purpose of the present paper, it is irrelevant whether the rapid increase of the density of states is due to the existence of extra dimensions, due to "stringy effects" becoming relevant at lower energies or to something else.

Let us now assume that we have a rapidly increasing level density, which has approximately the same form as in string theories. We choose:

$$d\left(m_n\right) \sim \left(\frac{m_n^2}{m_0^2}\right)^{-\alpha} \exp\left(\frac{m_n^2}{m_0^2}\right)^{\rho}, \tag{1}$$

see for instance [21]. Here, m_n stands for the mass of the resonance at

the n^{th} level of a string model, m_0 is the characteristic energy scale of the model. A numerical prefactor has been omitted: it is model dependent and it it does not play a significant role in what follows. Similarly, the exponent α is model dependent. We found, however that varying the exponent does not affect the results significantly; for the sake of definiteness, we settled with $\alpha = 2$.

We quoted an expression corresponding to the asymptotic form of the level density in a string model. Due to the fact that at this stage we want to understand the qualitative behavior of the neutrino cross section in the scenario just outlined, the present expression should be adequate. In most string models, the value of the exponent ρ is $1/2$. Nevertheless, we wanted to explore more rapidly rising level densities as well; all the estimates were carried out with $\rho = 1/2$ and $\rho = 1$.

In order to estimate the neutrino–quark cross section, we take into account the contribution of the resonances to the imaginary part of the forward neutrino–quark amplitude. The invariant Breit-Wigner formula is used for a single resonance. In this way, we get the contribution of a single resonance of mass squared s_n at level n of a string model:

$$ImB_n = \frac{1}{\pi} \frac{s \Gamma_n s_n^{1/2}}{(\hat{s} - s_n)^2 + \Gamma_n^2 s_n} \tag{2}$$

In this equation $\hat{s}$ stands for the CM energy squared of the neutrino-quark system.

The total cross section due to the "new physics" is then obtained by multiplying eq. (2) with the level density and summing over the levels. We assume that $\Gamma_n s_n^{1/2}$ grows linearly with s_n. *viz.* $\Gamma_n s_n^{1/2} = \gamma_0 s_n$ and that the resonances lie on a linear Regge trajectory, $s_n = s_0(n-1)$. Further, we identify $s_0 = m_0^2$ in eq. (1). (In a string model, both quantities are related to the string tension; there may be prefactors of order unity which we ignore here.) The linear growth of the width can be understood in intuitive terms by noticing that a resonance preferentially decays into channels with the largest phase space available. The number of such channels is roughly proportional to the mass of the resonance, hence $\Gamma_n s_n^{1/2} \propto s_n$. (It is known, for instance that the total widths of baryon resonances grow approximately linearly with their masses *cf.* [22].)

We now notice that γ_0 is likely to be rather small, *cf.* [11]. Hence, the quantity ImB_n may be replaced by a δ-function and the summation over the levels by an integration. In this approximation the neutrino–quark cross

section becomes:

$$\hat{\sigma} \approx \frac{1}{s_0} d\left(\hat{s}\right) \tag{3}$$

Finally, one has to integrate eq. (3) over the momentum distribution of quarks within the target nucleon, using the relation $\hat{s} = xs$. Little or nothing is known about the evolution of structure functions into a region of rapidly increasing level density. In order to get some idea, we chose a "generic" structure function, of the form:

$$S(x) = Ax^{-\gamma}\left(1-x\right)^{\delta} \tag{4}$$

Inspired by the Duke–Owens parametrization of the structure functions, *cf.* [23] we chose $A = 2, \gamma = 0.6, \delta = 3.5$. None of the results turned out to be very sensitive to the precise choice of these parameters. The final expression of the neutrino-nucleon cross section is given by:

$$\sigma = \int_{x_0}^{1} dx S(x) \hat{\sigma}(xs) \tag{5}$$

The infrared cutoff was chosen as $x_0 = s_0/s$; in this way, the "new physics" begins to manifest itself for $s \gtrsim s_0$. In order to test the scheme developed, let us assume that we want the cross section to grow to $\sigma \approx 1/\Lambda_{QCD}^2$ around the GZK cutoff, say, $s = 2 \times 10^4 \mathrm{TeV}^2$. (This corresponds to a laboratory energy, $E_L \approx 10^{19}$eV.) On taking $\Lambda_{QCD} = 200$MeV, one gets $s_0 \approx 3$TeV for $\rho = 1/2$ in eq. (1) and $s_0 \approx 30$TeV for $\rho = 1$, respectively. In the following figures we display the cross sections resulting from eq. (5) taking $\rho = 1/2$ and $\rho = 1$ in eq. (1).

One sees that, for all practical purposes, the neutrino-nucleon cross section is dominated by the exponential growth of the level density; hence, the choice of the exponents in eq. (1) and eq. (4) is not critical.

The cross sections rise very rapidly and, as it is always the case with tree amplitudes containing high spin particles in intermediate states, it will violate unitarity at a certain energy. The value of that energy is, however, difficult to determine at this stage. This is due to the fact that, like in any string model, at level n, resonances with spins $0 \leq s \leq n$ are exchanged. (In some models the lower limit may be 1 and the upper limit $n \pm 1$, but this fact does not affect the general situation.) As a consequence, the usual semiclassical estimates of the unitarity limits of cross sections are not directly applicable. Only by calculating loop corrections will one be able to estimate the unitarity limits.

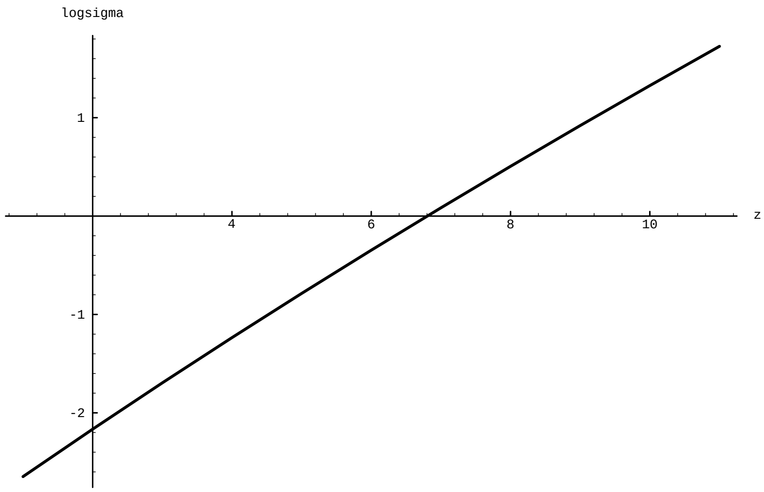

Figure 1: Decimal logarithm of the $\nu - N$ cross section measured in millibarns; $\rho = 1/2$, $s_0^{1/2} = 3\mathrm{TeV}$; $z = s/s_0$

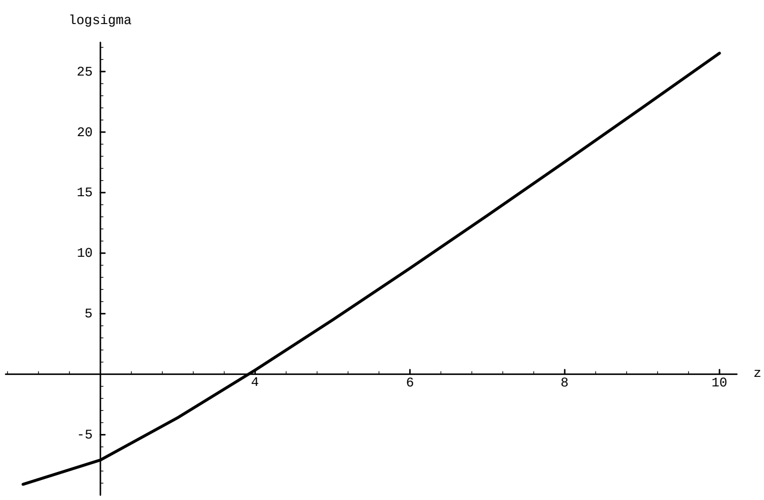

Figure 2: Same as Fig. 1, but with $\rho = 1$, $s_0^{1/2} = 30\mathrm{TeV}$

It is to be emphasized that *individual partial waves* need not violate the unitarity bounds. In order to illustrate this point, consider a simple Veneziano amplitude,

$$A(s,t) = \mathcal{B}\left(1-\alpha(s), 1-\alpha(t)\right), \tag{6}$$

where $\mathcal{B}(x,y)$ is the Euler beta function and α is a linear Regge trajectory. This amplitude has a level density asymptotically described by eq. (1) with $\rho = 1/2$. By using Stirling's formula, one obtains an asymptotic estimate for the partial wave amplitudes. In particular, one finds for the S wave amplitude:

$$A_0(s) \sim \left(2\pi\alpha(s)\right)^{-1/2} \quad (\alpha(s) \gg 1)\,. \tag{7}$$

Clearly, the bound given by $|\exp(i\delta_0)\sin\delta_0| \leq 1$ is satisfied asymptotically. (In order to arrive at the result given by eq. (7), one has to avoid the poles of the Γ function, *e.g.* by sending s to infinity along a ray in the complex plane, $s = |s|\exp i\phi, \quad \phi \neq 0$.)

Are the precision results of the Standard Model affected? This is an important question: the effects of the "new physics" are expected to be observable even below the characteristic energy scale, *cf.* Goldberg and Weiler, ref. [24]. As in discussing the unitary bounds before, we are unable to estimate loop effects at this stage. However, one can estimate final state interactions due to the "new physics". We use the scattering length approximation to the l^{th} partial wave,

$$k^{(2l+1)}\cot\delta_l \approx \frac{1}{a_l} \tag{8}$$

This should be reasonably accurate: typical CMS energies at precision measurements carried out at LEP and SLAC are of the order of 100GeV, whereas the characteristic energy of the "new physics" is on the TeV scale. A reasonable estimate for the scattering length is $1/a \approx \sqrt{s_0}$. The effect of the final state interactions is given by the formula, *cf.* [25]:

$$w_l \approx w_l^{(0)}\left(1 + \left(k^{(2l+1)}a_l\right)^2\right), \tag{9}$$

In eq. (9), w is a transition probability in a given partial wave l, (for instance, $Z \to \nu\bar{\nu}$) and w^0 is the same transition probability calculated ignoring the final state interaction. On using the estimate $\sqrt{s_0} \approx 3$TeV, one finds that with $k \approx 45$GeV (half of the mass of the Z), the S wave final state "enhancement factor" differs from unity by about 10^{-4} or so. This is to be compared with a typical error of a precision measurement (for instance, the

decay width of a weak gauge boson) which is of the order of 0.1%, [26]. The effect is, of course, much smaller in higher partial waves or for a higher characteristic energy, *cf.* eq. (9).

The scenario outlined here has some observable consequences. They, in turn, depend on whether the characteristic energy is of a few TeV or a few tens of TeV. For a characteristic energy of a a few TeV,

- there should be measurable deviations from standard model predictions in the decay rates of weak gauge bosons if the accuracy of the measurements can be increased by about an order of magnitude,
- one expects spectacular phenomena of new particle production and/or strong violation of Feynman scaling due to the rapid rise of the level density at LHC. Due to the precocious unification, one expects a copious production of leptons as well as hadrons.

For a high characteristic energy ($\simeq 30$TeV), this phenomenon may not take place at the LHC. However, new phenomena will be observed at nonaccelerator experiments, such as neutrino telescopes and at orbiting detectors (OWL, Airwatch).

- It was pointed out that orbiting detectors should be sensitive to neutrino interactions, see [27]. If the cross section of neutrino interactions shows a rapid rise, there should be a corresponding rise of the impact parameter of the incident neutrino with respect to the center of the Earth at which showers generated by them can be observed. Accordingly, there should be a cutoff in the spectrum of upward going neutrinos observed in underground, under water or under ice detectors. The neutrino induced showers develop rapidly within the Earth and they degenerate by the time they would reach the detector. By contrast, the number of showers of grazing incidence should increase.
- This scenario resolves the "energy crisis" caused by some schemes purporting to explain the highest energy cosmic rays, such as [9]. It is difficult enough to construct mechanisms by means of which protons are accelerated to energies of the order of 10^{20}eV, *cf.* Norman *et. al.* , [28]. If one prefers the highest energy cosmic rays to be neutrinos originating from the decay of pions and kaons, one needs protons to be accelerated to energies a few orders of magnitude even higher: at such energies, several hundred light hadrons are produced and, on the average, the available primary energy is shared equally between them. By contrast, in a scenario involving precocious unification, once the

CMS energy of the accelerated protons in an active galactic nucleus or in a similar site of intense proton acceleration reaches the characteristic energy, neutrinos are produced at multiplicities comparable to hadrons. Hence, there is no need to postulate proton energies several orders of magnitude larger.

We thank L. Madansky and C. Norman for several conversations on the highest energy cosmic rays.

References

[1] M. Takeda *et. al.* , Phys. Rev. Lett. **81**,1163 (1981); S. Yoshida *et. al.* , Astropart. Phys. **3**, 105 (1995); N. Hayashida *et. al.* , Phys. Rev. Lett. **73**, 3491 (1994). (AGASA).

[2] D.J. Bird *et. al.* , Phys. Rev. Lett. **71**, 3401 (1993); Ap. J. **424**, 491 (1994); *ibid.* **441**, (1995). (Fly's Eye).

[3] G. Brooke *et. al.* , Proc. 19^{th} Internat'l Cosmic Ray Conf. (La Jolla) **2**, 150 (1985); A.A. Watson, J. Phys. G **17**, 733 (1995). (Haverah Park).

[4] N.N. Efimov *et. al.* , in Proc. ICRR Symposium on Astrophysical Aspects of the Most Energetic Cosmic Rays, ed. N. Nagano and F. Takahara. World Scientific, Singapore (1991). (Yakutsk).

[5] J. Szabelski, preprint astro-ph/9710191

[6] C.T. Hill and D.N. Schramm, Phys. Rev. **D31**, 564 (1985).

[7] Workshop on Observing Giant Cosmic Ray Air Showers from $> 10^{20}$eV Particles from Space. Editors: J.F. Krizmanic, J.F. Ormes and R.E. Streitmatter. AIP Conference Proceedings 433. American Institute of Physics, Woodbury, N.Y. 1998.

[8] G. Burdman, F. Halzen and R. Gandhi, Phys. Lett. **B417**, 107 (1998).

[9] G. Domokos and S. Nussinov, Phys. Letters **B187**, 372 (1987)

[10] G. Domokos and S. Kovesi–Domokos, Phys. Rev. **D38**, 2833 (1988)

[11] I. Antoniadis, S. Dimopoulos and G. Dvali, Nucl. Phys. **B516**, 70 (1997.

[12] N. Arkani–Hamed, S. Dimopoulos and G. Dvali, Phys. Lett. **B429**, 263 (1998).

[13] N. Arkani–Hamed, S. Dimopoulos and J. March–Russell, preprint hep-th/9809124 and references quoted there.

[14] K.R. Dienes, E. Dudas and and T. Ghergetta, Phys. Letters **B436**, 55 (1998).

[15] K.R. Dienes, E. Dudas and T. Ghergetta, hep-ph/9807522.

[16] K.R. Dienes, E. Dudas, T. Ghergetta and A. Riotto, preprint hep-ph/9809406 and references quoted there.

[17] C. Giunti, C.W. Kim and U.V. Lee, Mod. Phys. Lett. **A6**, 1475 (1991).

[18] U. Amaldi, W. de Boer and H. Fürstenauer, Phys. Lett. **B260**, 447 (1991).

[19] J. Ellis, S. Kelley and D.V. Nanopoulos, Phys. Lett. **B260**, 131 (1991).

[20] P. Langacker and M. Luo, Phys. Rev. **D44**, 817 (1991).

[21] See, *e.g.* M.B. Green, J.H. Schwarz, and E. Witten Superstring Theory, Cambridge University Press, Cambridge, 1987. Vol. 1, p.118.

[22] G. Domokos, S. Kovesi–Domokos and E. Schonberg, Phys. Rev. **D3**, 1184 (1971).

[23] V.D. Barger and R.J.N. Phillips, Collider Physics. Addison–Wesley, Redwood City, CA 1987. Ch. 5.

[24] H. Goldberg and T.J. Weiler, preprint hep-ph9810533.

[25] M.L. Goldberger and K.M. Watson, Scattering Theory. Wiley, New York, 1964. Sec. 9.3.

[26] Particle Data Group, http://pdg.lbl.gov

[27] G. Domokos and S. Kovesi–Domokos, hep-ph/98052221 and in ref. [7]

[28] C.A. Norman, D.B. Melrose and A. Achterberg, Ap.J. **454**, 60 (1995).